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MOMENTUM DEPENDENT VERTICES $\sigma\gamma\gamma$, $\sigma\rho\gamma$ and $\sigma\rho\rho$: THE NJL SCALAR HIDDEN BY CHIRAL SYMMETRY [1]

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ABSTRACT: We calculate the momentum dependence of three particle vertices $\sigma\gamma\gamma$, $\sigma\rho\gamma$ and $\sigma\rho\rho$ in the context of a Nambu Jona Lasinio type model. We show how they influence the processes $\gamma\gamma \to \sigma \to \pi\pi$, $\rho \to \gamma\sigma$ and $\gamma\gamma \to \rho\rho$ and how chiral symmetry shadows the presence of the σ.

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1 INTRODUCTION

Presently most of the properties of low energy hadron physics are described in terms of effective lagrangians which incorporate some of the important symmetries of QCD. The Nambu Jona Lasinio (NJL) type lagrangians are particularly popular due to the fact that they provide for an adequate description of low energy meson masses, their decays, scattering, etc. in terms of quarks [1]-[5]. One of the important predictions of the NJL model is the existence of a broad scalar excitation - the σ particle. Its mass is exactly equal to $2M_u$ (M_u is the constituent up-quark mass) in the chiral limit (current quark masses $m_u = m_d = 0$ or, equivalently, $m_\pi = 0$) of the 2 flavours NJL and it is a little bit smaller than $2M_u$ in a more general framework ($m_\pi \neq 0$, strange quark contribution included and $U(1)_A$ symmetry explicitly broken), which we will adopt. The experimental evidence of the σ resonance is however controversial. Since its non-existence would enter in contradiction with the NJL model, it is one of the purposes of the present article to show that the experimental results on the $\gamma\gamma \to \pi\pi$ and $\rho \to \gamma\pi\pi$ processes are by no means in conflict with the presence of the σ meson predicted by the NJL. Moreover, the σ meson may resolve some puzzles connected with the reactions $\gamma\gamma \to \rho\rho \to 4\pi$ and $pp \to pp\pi^0$, as discussed in the sequel.

The amplitude analysis [6] of the data [7]-[8] for the reaction $\gamma\gamma \to \pi\pi$ shows that the σ- contribution in the s-channel near the pole cannot be large. The same can be said for the decay $\rho^0 \to \gamma\sigma$ [9]: the branching ratio for $\rho^0 \to \gamma\pi^+\pi^-$ is in fact 1 %, but the absence of $\rho^0 \to \gamma\pi^0\pi^0$ clearly means that the two pions cannot be in a S-wave. It will be shown that the intermediate σ meson in reactions involving photons is hidden due to the chiral symmetry prediction $m_\sigma = 2M_u$.

The reaction $\gamma\gamma \to \rho\rho \to 4\pi$ shows a large enhancement of the $\rho^0\rho^0$ [10] intermediate channel as compared to the $\rho^+\rho^-$ [11] channel. Several models give partial answers to the problem. Within the MIT bag quark model, the candidates for this reaction appear as $2q - 2\bar{q}$ states with masses around 1.65 GeV [12]. Using the multiplets suggested in [12], but with a little lower masses (1.4 GeV), the $\rho^0\rho^0$ channel is quite well reproduced, while the so obtained $J^P = 2^+$ above 1.7 GeV and all the 0^+ contributions to the $\rho^+\rho^-$ channel are too small [13]. In the effective lagrangian approach [14], using π, η, η', ρ, ω, ϕ, $\epsilon(1300)$, $f_2(1270)$, and $a_2(1320)$, the K-matrix method is employed to unitarize the amplitude. Similar to the quark-inspired model [13], the $\rho^0\rho^0$ cross section is well reproduced, while the cross section obtained for $\rho^+\rho^-$ is too large by a factor of two.

We propose another line of investigation, namely the hypothesis that the reaction proceeds via a scalar exchange, because the two photons would be converted only in $\rho^0\rho^0$, due to zero charge exchange. We are encouraged to consider this mechanism by a simple qualitative estimate in which we adapt the Yukawa-type scalar exchange potential proposed by Machleidt et al.[15] to the $\rho\rho$ system by multiplying with a factor $(\frac{2}{3})^2$ since there are two quarks rather than three at each vertex. Choosing a reasonable hard core at .17 fm and solving a non-relativistic Schrödinger equation, we obtain at $E \approx 0$ a weakly bound or resonant state. Our calculation of the q^2 dependence of the form factors for $\sigma\rho\gamma$ and $\sigma\rho\rho$ shows that the existence of the σ-meson could resolve the puzzle offered. In this case the $\sigma\rho\gamma$ vertex contributes in a kinematical region far from the one where chiral symmetry predicts shadowing of the scalar.

To complete the picture of the scalar meson predicted by the NJL model, we mention

its relevance also in cases which do not involve photons and where no shadowing due to chiral symmetry is observed. One example is $\pi\pi$ scattering in some models [3], where the inclusion of the scalar is crucial for the description of the data on scattering lengths and phase shifts. The other example is the reaction $pp \to pp\pi^0$ where the hypothetical σ meson can also resolve a puzzle, namely the underestimate of the cross section [16]. As pointed out by Horowitz et al. [17], the inclusion of the σ exchange between a positive and a negative energy virtual state of the intermediate protons can in fact account for the discrepancy by a factor of 5 between experimental results and theoretical expectations.

2 THE MODEL AND FORM FACTORS

The following NJL lagrangian has been shown to provide for an adequate description of mesonic mass spectra and radiative decays [5,18]

$$
\begin{aligned}
\mathcal{L} = \quad & \bar{\psi}[\,i\partial\!\!\!/ - m_{curr} - \frac{ie}{2}(\lambda_3 + \frac{\lambda_8}{\sqrt{3}})A\!\!\!/]\psi \\
+ \quad & G_1[(\bar{\psi}\lambda_i\psi)^2 + (\bar{\psi}i\gamma_5\lambda_i\psi)^2] \\
+ \quad & G_2[(\bar{\psi}\lambda_a^c\gamma_\mu\psi)^2 + (\bar{\psi}\lambda_a^c\gamma_\mu\gamma_5\psi)^2] \\
+ \quad & K[\det\{\bar{\psi}(1+\gamma_5)\psi\} + \det\{\bar{\psi}(1-\gamma_5)\psi\}]
\end{aligned} \tag{1}
$$

which is an SU(3) generalization [2] of the NJL model [1], coupled minimally to the electromagnetic interaction. Here the flavor index i runs from 0 to 8 with $\lambda_0 = \sqrt{\frac{2}{3}}\mathbf{1}$. The λ_a^c are color matrices (a=1,...,8) and m_{curr} is the diagonal current quark mass matrix. The usual regularization scheme is used with a covariant four-momentum cutoff, $\Lambda = 1$ GeV. The parameters used are the same as in [5]: $G_1\Lambda^2 = 3.95$, $G_2\Lambda^2 = 5.43$, $K\Lambda^5 = 42$ and the quark current masses $m_u = m_d = 4$ MeV, $m_s = 115$ MeV. The constituent quark masses M_i for the given model parameters are $M_u = M_d = 390$MeV and $M_s = 536$MeV. In the present calculation we are interested in vertices of the type $\sigma VV'$, where V, V' stand for vector mesons or photons and σ is the scalar composed of $\lambda_0 - \lambda_8$ flavor admixtures. These vertices are obtained by evaluating the Feynman diagram shown in fig.1 , where the triangle is composed of quark lines. The full electromagnetic vertex, denoted by $\Gamma_\mu^{(\gamma)}$, is given by [18]

$$
\begin{aligned}
\Gamma_\mu^{(\gamma)}(q^2) = \quad & i\gamma_\mu \textstyle\sum_{i=0,3,8} \lambda_i \mathcal{G}_i^{(\gamma)}(q^2) \\
= \quad & i\frac{e}{2}\gamma_\mu[(\lambda_3 + \frac{\lambda_8}{\sqrt{3}}) + \textstyle\sum_{i,j=0,3,8} \lambda_i (J_{3j}(q^2) + J_{8j}(q^2))\mathcal{M}_{ji}(q^2)]
\end{aligned} \tag{2}
$$

The quantities J_{3j} and J_{8j} are the fundamental quark bubbles which couple the bare electromagnetic vertex to the vector particles ρ, ω and ϕ. These particles are generated through the coupling to the strong quartic $q\bar{q}$ interaction, represented by the 3×3 matrices $\mathcal{M}_{ji}$, where i and j are flavor indices. The electromagnetic Ward identity for this vertex is fulfilled. In the present case, with the SU(2) flavor symmetry preserved in the up and down sector, the ρ - meson decouples from the ω - ϕ excitations and $\mathcal{M}_{ji}$ reduces to a number in the case of the ρ - meson and to a 2×2 submatrix which describes ω - ϕ mixing.

The off-shell $\rho q\bar{q}$ coupling $\mathcal{G}_3^{(\rho)}(q^2)$ can be extracted from the $q\bar{q}$ scattering matrix $\mathcal{M}_{33}$ by separating the pole contribution and is given by

$$\Gamma_\mu^{(\rho)}(q^2) = i\gamma_\mu\lambda_3\mathcal{G}_3^{(\rho)}(q^2) = i\gamma_\mu\lambda_3\sqrt{\frac{16}{9}G_2\frac{q^2-m_\rho^2}{1-\frac{16}{9}G_2J_V(q^2)}} \tag{3}$$

where J_V is the quark bubble which appears in the Bethe Salpeter solution to one-loop order in the vector-isovector channel and m_ρ the ρ - meson mass obtained at the corresponding pole. The coupling is momentum dependent and contains all the information about the $q\bar{q}$ structure of the ρ meson. In a similar way the $\sigma q\bar{q}$ couplings $\mathcal{G}_0^{(\sigma)}$ and $\mathcal{G}_8^{(\sigma)}$ are extracted from the corresponding $q\bar{q}$ scattering matrix; for details see e.g. [4]. Following Feynman rules we obtain for the $\sigma VV'$ amplitude, with $V, V' = \gamma$ or ρ

$$A_{\sigma VV'} = \epsilon_{3b}^\mu\epsilon_{2a}^\nu \sum_{n=0,8}\sum_{i=u,d,s} M_i\mathcal{G}_n^{(\sigma)}(q_1^2)X_{i,n}^{VV'}(q_2^2,q_3^2)T_{\mu\nu}^i(q_2,q_3) \tag{4}$$

where $q_3 = q_1 + q_2$ and

$$T_i^{\mu\nu} = 6i\int_\Lambda \frac{d^4k}{(2\pi)^4}\frac{Tr_D\gamma^\mu(\not{k}+M_i)\gamma^\nu(\not{k}-\not{q}_2+M_i)\mathbf{1}(\not{k}-\not{q}_3+M_i)}{[k^2-M_i^2][(k-q_2)^2-M_i^2][(k-q_3)^2-M_i^2]} \tag{5}$$

The factors $X_{i,n}^{VV'}$ are the result of the flavor trace in the triangle and involve the momentum dependent couplings of the photon or ρ to the quarks. The factor 6 results from exchange and from taking the trace over color. We have explicitly for V and V' being photons

$$X_{i,0}^{VV'} = \sqrt{\frac{2}{3}}\sum_{n,m=0,3,8}\mathcal{G}_n^{(\gamma)}(q_2^2)\mathcal{G}_m^{(\gamma)}(q_3^2)c_{nm}^i \tag{6}$$

with $c_{nm}^i = c_{mn}^i$ where $c_{00}^u = \sqrt{2}c_{08}^u = 2c_{88}^u = \frac{2}{3}c_{33}^u = \sqrt{2}c_{38}^u = c_{03}^u = \frac{2}{3}$, all $c_{nm}^d = c_{nm}^u$ except $c_{n3}^d = -c_{n3}^u$ $(n \neq 3)$, and $c_{00}^s = -\frac{1}{\sqrt{2}}c_{08}^s = \frac{1}{2}c_{88}^s = \frac{2}{3}$. The coefficients c_{n3}^s equal zero. One has furthermore $X_{i,0}^{VV'} = \frac{1}{\sqrt{2}}X_{i,8}^{VV'}$; $(i = u, d)$ and $X_{s,0}^{VV'} = -\frac{1}{2}X_{s,8}^{VV'}$. In the case of one or

both V and V' being a ρ - meson one has to substitute correspondingly $\mathcal{G}_n^{(\gamma)}$ by the ρ-$q\bar{q}$ coupling $\mathcal{G}_3^{(\rho)}$ and there is obviously no triangle with strange quarks involved in (4).

The expression (4) for the vertex function $A_{\sigma VV'}$ can be written in a more compact form

$$A_{\sigma VV'} = \epsilon_{3b}^\mu\epsilon_{2a}^\nu T_{\mu\nu}(q_2,q_3) \tag{7}$$

In order to evaluate $T_{\mu\nu}$ we decompose in gauge invariant tensors [19] as

$$T_{\mu\nu} = F_{TT}G_{\mu\nu} + F_{LL}L_{\mu\nu} \tag{8}$$

with orthogonal projection operators

$$G_{\mu\nu} = g_{\mu\nu} + (q_3^2q_{2\mu}q_{2\nu} + q_2^2q_{3\mu}q_{3\nu} - q_2q_3(q_{2\mu}q_{3\nu} + q_{3\mu}q_{2\nu}))/Y \tag{9}$$

$$L_{\mu\nu} = q_2 q_3 \frac{1}{Y} Q_{3\mu} Q_{2\nu} \tag{10}$$

where $Q_{3\mu} = q_{3\mu} - (q_3^2/q_3 q_2) q_{2\mu}$, $Q_{2\nu} = q_{2\nu} - (q_2^2/q_3 q_2) q_{3\nu}$ and $Y = (q_2 q_3)^2 - q_2^2 q_3^2$.

From eqs. (4), (7) and (8) it is easy to see that

$$F_{TT(LL)}(q_1^2, q_2^2, q_3^2) = \sum_{n=0,8} \sum_{i=u,d,s} M_i \mathcal{G}_n^{(\sigma)}(q_1^2) X_{i,n}^{VV'}(q_2^2, q_3^2) F^i_{TT(LL)}(q_1^2, q_2^2, q_3^2). \tag{11}$$

In order to evaluate the transverse F_{TT} and longitudinal F_{LL} form factors, we need first the triangle contributions of each flavor separately, which are obtained with the contractions

$$F^i_{TT} = \frac{1}{2} G_{\mu\nu} T_i^{\mu\nu}$$

and

$$F^i_{LL} = \frac{(q_3 q_2)^2}{q_3^2 q_2^2} L_{\mu\nu} T_i^{\mu\nu}.$$

The amplitude $T_{\mu\nu}$ from (7) is calculated with the following prescription (which is in a sense a definition of the cutoff): first, the reduction methods of ref.[20] are used; the integral is treated at this stage as without cutoff, allowing in this way translation of the integration variable. As a result, $T_{\mu\nu}$ is expressed as a linear combination of the basic integrals of the form (13) below and in this way the μ, ν dependence of the integration variable removed. Second, a sharp Euclidean cutoff is introduced for the scalar integrals I_2 and I_3 (13). All the cutoff dependence of $T_{\mu\nu}$ thus resides in these basic integrals. The result could be in principle gauge dependent, due to the introduction of a sharp Euclidean cutoff. A necessary condition for gauge independence of $T_{\mu\nu}$ is that the Ward identities $q_3^\mu T_{\mu\nu} = q_2^\nu T_{\mu\nu} = 0$ are fulfilled. We can show, that these Ward identities are indeed exactly fulfilled, regardless of the cutoff Λ, with no need to specify what I_2 and I_3 are (the coefficients in front of I_2 and I_3 in the expressions $q_3^\mu T_{\mu\nu}$ and $q_2^\nu T_{\mu\nu}$ vanish exactly). In this way we check that the above explained recipe for handling the initial expression (7) is meaningful.

The results for the form factors can be cast in the form

$$F^i_{TT} = \frac{1}{2} [\mathcal{A}_i - \frac{q_2 q_3}{Y} \mathcal{H}_i] \tag{12}$$

and

$$F^i_{LL} = \frac{q_2 q_3}{Y} \mathcal{H}_i$$

with

$$\mathcal{A}_i = g_{\mu\nu} T_i^{\mu\nu} =$$
$$4(I_3^i \times (4 q_2 q_3 + 4 M_i^2 - q_3^2 - q_2^2) + I_2^i(q_2^2) + I_2^i(q_3^2) - 2 I_2^i((q_2 - q_3)^2)))$$

and

$$\mathcal{H}_i = q_{2\mu} q_{3\nu} T_i^{\mu\nu} =$$
$$4((q_2 q_3 - q_2^2)(q_2 q_3 I_3^i + I_2^i(q_2^2)) + (q_2 q_3 - q_3^2)(q_2 q_3 I_3^i + I_2^i(q_3^2)) + (q_2 - q_3)^2 I_2^i((q_2 - q_3)^2)))$$

with the standard integrals

$$I_2^i(q^2) = i \int_\Lambda \frac{d^4k}{(2\pi)^4} \frac{1}{[k^2 - M_i^2][(k+q)^2 - M_i^2]} \tag{13}$$

and

$$I_3^i = I_3^i(q_2^2, q_3^2, q_2 q_3) =$$
$$i \int_\Lambda \frac{d^4k}{(2\pi)^4} \frac{1}{[k^2 - M_i^2][(k-q_2)^2 - M_i^2][(k-q_3)^2 - M_i^2]}.$$

3 RESULTS AND DISCUSSION

With the choice of parameters of ref.[5] we obtain the masses $m_\sigma = 775$ MeV and $m_\rho = 775$ MeV. Let's start to discuss the results for the $\sigma\gamma\gamma$ vertex, which may serve as a guide to extract information about the existence of a broad resonant state in the reaction $\gamma\gamma \to 2\pi$. The interesting kinematical region is with two on-shell photons and a time-like scalar near its on-shell value. The form factor F_{TT} for this case is displayed in fig.2. The longitudinal form factor F_{LL} is for on-shell photons of course irrelevant, because they cannot be longitudinal. An interesting feature can be seen for the transversal form factor F_{TT}: it is very small for an on-shell σ-meson. It can be proved analytically from eq. (12) and the definitions for $\mathcal{A}_i$ and $\mathcal{H}_i$ that $F_{TT}^i = 2(4M_i^2 - q_1^2)I_3^i$ for real photons. In the exact chiral $U(3)_L \times U(3)_R$ symmetry limit the constituent masses of all the u, d and s quarks are equal and one obtains the relation $m_\sigma = 2M_u$. In this limit F_{TT} for $\sigma\gamma\gamma$ vanishes exactly for on-shell photons and σ. In our case the chiral $U(3)_L \times U(3)_R$ symmetry is explicitly broken by $m_{curr} \neq 0$ and $K \neq 0$, so that the form factor F_{TT} is not exactly zero but close to it. This means, that the contribution of the σ to the process $\gamma\gamma \to \pi\pi$ is very small close to the σ mass. We will discuss in the sequel that the zero of the relevant on-shell form factor for $\sigma\gamma\gamma$ in the chiral limit is the underlying physics which forbids the observation of a scalar resonance in processes such as $\gamma\gamma \to \pi\pi$ and $\rho \to \gamma(\pi\pi)_{S-wave}$. It will explain how a scalar excitation can exist, but not contribute to certain processes. To obtain the complete NJL prediction to the $\gamma\gamma \to \pi\pi$ reaction via scalar exchange, one needs to couple the $\sigma\gamma\gamma$ vertex to the two pions in the exit channel, taking in account the decay width of the scalar into pions. This corresponds to the 'resonant' contribution. For the model parameters used the coupling of the NJL scalar to two pions yields a broad resonance width $\Gamma_{\sigma\pi\pi} = 435$MeV at the σ mass. A simple estimate of the cross section for this process can be obtained assuming a Breit-Wigner type resonance, with the width as a function of the square momentum of the σ predicted in NJL. The contribution of the σ meson should not be large, since it would otherwise enlarge too much the experimentally small cross–section in the $\pi^0\pi^0$ channel [8]. We found that within the NJL and with the Breit-Wigner form for the σ propagator the resonant contribution is very small near the

scalar mass (≈ 0.3 nb at $W_{\gamma\gamma} = 0.775$ GeV), but is increasing with decreasing energy (≈ 10 nb at 0.5 GeV) until it becomes too large (≈ 50 nb at 0.32 GeV). Nevertheless, near the sigma mass, the resonant scalar contribution is shadowed by chiral symmetry, so that the existence of the σ is not in contradiction with the data in this kinematical region.

A comment is in order concerning a complete calculation of the $\gamma\gamma \to \pi^0\pi^0$ reaction within the NJL model: processes other than the scalar exchange considered presently are involved to the same order, like a quark box diagram with two photons and two pions attached to the vertices, and the exchange of vector mesons in the u and t - channels. These processes can be denoted as interference or non-resonant terms. However, for the purposes of showing that the predicted broad NJL scalar is compatible with the small cross sections observed below the nominal σ mass in the reaction $\gamma\gamma \to \pi^0\pi^0$, it is enough to consider the "resonant" term for the following reasons. We calculate the vector meson exchange for $cos(\theta) = 0$ and find it to be negligible in the kinematical region considered. We expect the box diagram (which we have not calculated) to lower the cross section, since the σ exchange graph and the box diagram cancel in the chiral limit and for photon frequencies $\omega_i \to 0$, see [21]. This fact is especially welcome near threshold, where the resonant contribution predicted by the NJL is too large. A complete calculation is rather involved and relegated to a later work. In our opinion it will not alter the conclusion that chiral symmetry is at work producing a small $\sigma\gamma\gamma$ amplitude, which vanishes in the chiral limit (due to $m_\sigma = 2M_u$) and is responsible for the small resonant scalar contribution observed around the nominal σ mass.

The screening of the σ contribution is present also in the $\rho \to \gamma(\pi\pi)_{(S-wave)}$ decay. The intermediate σ must be somewhat off-shell because of phase space. From fig. 3 it is in fact possible to see that the F_{TT} of the $\sigma\rho\gamma$ vertex is small whenever the scalar and the photon are nearly on-shell. In the case of one on-shell photon the transverse form factors reduce to

$$F^i_{TT} = 2((4M_i^2 - q_1^2 + q_2^2)I_3^i - \frac{2q_2^2}{q_2^2 - q_1^2}(I_2^i(q_2^2) - I_2^i(q_1^2)))$$

as can be obtained again from (12). One observes in fig. 3 that F_{TT} is almost insensitive to the variations in q_2^2 of the ρ-meson at the on-shell values of the scalar and photon. Since F_{TT} is close to 0 for on-shell σ and $q_2^2 = 0$ (similar situation as the one discussed above for $\sigma\gamma\gamma$), it remains close to 0 also for a nearly on-shell ρ-meson. In the same figure we notice a steep increase of the form factors as functions of increasing absolute values of the squared space-like momenta of the scalar. The exchange of a *space-like* σ in the diagram with two $\sigma\rho\gamma$ vertices can give a sizeable contribution to the process $\gamma\gamma \to \rho^0\rho^0$, see fig.3.

Significant changes are also seen in the $\sigma\rho\rho$ form factors as shown in figs. 4a and 4b. These form factors are also relevant for the $\rho\rho$ interaction via scalar exchange. Since the form factors drop rather rapidly at the on-shell values of the ρ -mesons near $q_1^2 \simeq 0$ as functions of the q_1^2 of the σ meson , the Fourier transform of $F_{TT}^2/(q^2 - m_\sigma^2)$ will be shallower and of longer range than the corresponding Yukawa potential. However such effective potentials are dependent on the off-shell mass of the ρ-mesons (see figs.4); for lower q^2 of the ρ the potential becomes deeper and of shorter range. This might explain, why the enhancement remains large even below the nominal $\rho^0\rho^0$ threshold. Namely, "lighter" ρ-mesons (off-shell ρ-mesons) would feel a deeper potential and "resonate" at

the lower energy. This effect might make the resonance "follow" into the region below 1.5 GeV, rendering it broader and noticeable below 1.5 GeV.

4 CONCLUSIONS

To summarize, we have shown that the $\simeq$ 400MeV wide scalar resonance predicted by the NJL model is a good candidate to describe several radiative and strong interaction phenomena. In particular we have discussed the remarkable fact that the presence of the σ excitation in the reactons $\gamma\gamma \to \pi\pi$ and $\rho \to \gamma(\pi\pi)_{(S-wave)}$ is hidden due to chiral symmetry. This results are a beautiful illustaration of chiral symmetry at work at the scale of the σ mass. In the case of the reaction $\gamma\gamma \to \rho\rho \to 4\pi$, where a large enhancement of the $\rho^0\rho^0$ intermediate channel as compared to $\rho^+\rho^-$ is observed, we have proposed a picture in which the reaction proceeds via a scalar exchange. We have calculated the vertices of relevance for this process, $\sigma\rho\gamma$ and $\sigma\rho\rho$ and discussed the reasons for the enhancement. The vertices involving the σ meson calculated here can be used in the evaluation of cross sections. Work in this direction is in progress.

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FIGURE CAPTIONS

fig.1 Diagram to calculate the vertex $\sigma VV'$; V,V'=photon or vector meson (wavy lines) with momenta q_2 and q_3. The σ - mode has momentum q_1. The large circles with crosses and the full circle are to remind that the vertices are momentum dependent.

fig.2 Transverse F_{TT} form factor for the vertex $\sigma\gamma\gamma$ with both photons on-shell, as function of q_1^2 of the scalar.

fig.3 F_{TT} for $\sigma\rho\gamma$ with on shell photon. The abscissa is the off-shell mass squared of the scalar. Curves: a) with ρ-meson on shell; b) $q_2^2 = .5$ GeV^2; c) $q_2^2 = .3$ GeV^2.

fig.4a F_{TT} for $\rho\rho\sigma$. The abscissa is the off-shell mass squared of the scalar. Curves: a) off-shell masses of the ρ - mesons $q_2^2 = q_3^2 = .6$ GeV^2; b) $q_2^2 = .5$ GeV^2, $q_3^2 = .6$ GeV^2; c) $q_2^2 = q_3^2 = .5$ GeV^2

fig.4b Same as in 4a but for F_{LL}.